

Mass Degeneracy of the Higgsinos

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Abstract

The search for charginos and neutralinos at LEP2 can become problematic if these particles are almost mass degenerate with the lightest neutralino. Unfortunately this is the case in the region where these particles are higgsino-like. We show that, in this region, radiative corrections to the higgsino mass splittings can be as large as the tree-level values, if the mixing between the two stop states is large. We also show that the degree of degeneracy of the higgsinos substantially increases if a large phase is present in the higgsino mass term μ .

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The search for charginos ($\tilde{\chi}^+$) at LEP2 is one of the most promising ways of discovering low-energy supersymmetry. If the $\tilde{\chi}^+$ decays into the lightest neutralino ($\tilde{\chi}^0$) and a virtual W^+ , it can be discovered at LEP2 (with a $\int \mathcal{L} = 500 \text{ pb}^{-1}$) whenever its production cross section is larger than about 0.1–0.3 pb and $m_{\tilde{\chi}^0}$ is within the range $m_{\tilde{\chi}^0} \gtrsim 20 \text{ GeV}$ and $m_{\tilde{\chi}^+} - m_{\tilde{\chi}^0} \gtrsim 5\text{--}10 \text{ GeV}$ [1]. Therefore, the chargino can be discovered almost up to the LEP2 kinematical limit, unless one of the following three conditions occurs:

i) The sneutrino ($\tilde{\nu}$) is light and the chargino is mainly gaugino-like. In this case the $\tilde{\nu}$ t -channel exchange interferes destructively with the gauge-boson exchange and can reduce the chargino production cross section below the minimum values required for observability, 0.1–0.3 pb. However, in a large fraction of the parameter space where this effect is important, the two-body decay mode $\tilde{\chi}^+ \rightarrow \tilde{\nu} l^+$ is kinematically allowed and dominates over the conventional three-body decays $\tilde{\chi}^+ \rightarrow \tilde{\chi}^0 l^+ \nu$, $\tilde{\chi}^0 \bar{q} q$. The resulting signal, quite similar to the one caused by slepton pair production, allows the chargino search for much smaller production cross section, possibly as small as 20–60 fb [1].

ii) The $\tilde{\chi}^0$ is very light ($m_{\tilde{\chi}^0} \lesssim 20 \text{ GeV}$). In this case the $\tilde{\chi}^+$ -detection efficiency diminishes, as the decrease in missing invariant mass makes the signal more similar to the $W^+ W^-$ background. Such a light $\tilde{\chi}^0$ is allowed by LEP1 data, if the weak-gaugino mass (M) is small¹, $M \ll M_W$. It is however ruled out by gluino (\tilde{g}) searches at the Tevatron [3] in models which assume gaugino mass unification, $M = \alpha M_{\tilde{g}} / (\alpha_s \sin^2 \theta_W)$. It should also be mentioned that no study has attempted to optimize the analysis in the low- $m_{\tilde{\chi}^0}$ region, while specially designed experimental cuts could improve the chargino detection efficiency.

iii) $\Delta_+ \equiv m_{\tilde{\chi}^+} - m_{\tilde{\chi}^0}$ is small ($\Delta_+ \lesssim 5\text{--}10 \text{ GeV}$) and the $\tilde{\chi}^+$ detection is problematic because of the lack of energy of the visible decay products [1].

In this letter, we concentrate on case (*iii*). Small values of Δ_+ occur when the gaugino masses are much larger than M_W . In this region, the lightest chargino and the two lightest neutralinos ($\tilde{\chi}^0, \tilde{\chi}^{0'}$) are mainly higgsino-like and are nearly degenerate² with mass $\sim \mu$. In ref. [5] it has been suggested that this parameter region ($M \gg M_W$ and $\mu \sim M_W$), although problematic for chargino searches, can nevertheless be covered by neutralino searches. Indeed, in this case the $\tilde{\chi}^0\text{--}\tilde{\chi}^{0'}$ production cross section at LEP2 is large ($\sim \text{pb}$)

¹The parameter region where M and the higgsino mass (μ) are both small has been recently studied in ref. [2].

²It has been recently shown [4] that single-photon tagging cannot be used to observe charginos in the higgsino region under consideration. Ref. [4] has also considered a gaugino-like region where Δ_+ can become small in models without gaugino mass unification.

and the mass difference $\Delta_0 \equiv m_{\tilde{\chi}^{0'}} - m_{\tilde{\chi}^0}$ is always greater than Δ_+ , allowing a better identification of the visible decay products than in the chargino case. It is also worth recalling that the parameter region where $M \gg M_W$ and $\mu \sim M_W$ has a special interest as a light higgsino-like chargino (together with a light stop) can increase the Standard Model prediction for R_b [6].

Here we point out that, for $M \gg M_W$ and $\mu \sim M_W$, Δ_0 and Δ_+ receive one-loop corrections proportional to m_t^3 , which can be as large as their tree-level values. We will also show that in the small $\tan \beta$ region, the tree-level values of Δ_0 and Δ_+ can be reduced if the μ parameter has a non-trivial phase. Finally, we will comment on the implications for chargino and neutralino searches at LEP2.

First let us recall that, in the higgsino region under consideration ($M \gg M_W$, $\mu \sim M_W$), the tree-level values of Δ_0 and Δ_+ are well approximated by a $1/M$ expansion³

$$\begin{aligned}\Delta_0 &= 2a \frac{M_W^2}{M} + 2b \sin 2\beta \frac{\mu M_W^2}{M^2} + \mathcal{O}(1/M^3), \\ \Delta_+ &= [a + (a-1)\varepsilon \sin 2\beta] \frac{M_W^2}{M} + [(b-1) + b\varepsilon \sin 2\beta] \frac{|\mu| M_W^2}{M^2} \\ &\quad - \frac{a^2}{2} \cos^2 2\beta \frac{M_W^4}{|\mu| M^2} + \mathcal{O}(1/M^3),\end{aligned}\tag{1}$$

where

$$\begin{aligned}a &\equiv \frac{4}{5} + \frac{1}{2} \left(\frac{M}{M'} \tan^2 \theta_W - \frac{3}{5} \right), \\ b &\equiv \frac{1}{2} \left(1 + \frac{9}{25 \tan^2 \theta_W} \right) + \frac{1}{2} \left(\frac{M}{M'} \tan^2 \theta_W - \frac{3}{5} \right) \left(\frac{M}{M'} + \frac{3}{5 \tan^2 \theta_W} \right).\end{aligned}\tag{2}$$

Here $\tan \beta$ is the ratio of the two Higgs vacuum expectation values, $\varepsilon = \mu/|\mu|$ and M' is the hypercharge gaugino mass. The expansions in eq. (1) break down when $\mu \rightarrow 0$ but this is not relevant here since, for $M \gg M_W$, LEP1 data require $\mu > M_Z/2$.

By comparing the leading $1/M$ terms in eq. (1), one finds that $\Delta_0 > \Delta_+ > 0$ for any (positive) value of M' . Assuming the gaugino mass unification condition $M' = \frac{5}{3} \tan^2 \theta_W M$, we obtain:

$$\Delta_+ = \frac{1}{2} \left(1 - \frac{\varepsilon \sin 2\beta}{4} \right) \Delta_0 + \mathcal{O}(1/M^2).\tag{3}$$

³ Here we take M to be real and positive and μ to be real following the sign conventions of ref. [7]. A complex μ will be considered later.

Notice also that the critical region for chargino searches ($\Delta_+ \lesssim 10$ GeV) occurs when $M \gtrsim 400\text{--}600$ GeV, depending on $\tan\beta$. For the neutralinos, however, we have $\Delta_0 \lesssim 10$ GeV for $M \gtrsim 1$ TeV.

General expressions for the one-loop corrections to chargino and neutralino masses are given in ref. [8]. To obtain simple analytical formulae we have computed the radiative corrections in the limit $M \rightarrow \infty$. The only contributions arise from heavy quark–squark loops and $\gamma(Z)$ –higgsino loops⁴:

$$\delta\Delta_0 = 2G_t^2 m_t \sin 2\theta_t \sum_{i=1,2} (-1)^{i+1} B_0(\mu^2, m_t^2, m_{\tilde{t}_i}^2) + (t \leftrightarrow b), \quad (4)$$

$$\begin{aligned} \delta\Delta_+ = & \frac{\delta\Delta_0}{2} - \sum_{i=1,2} \left[G_t G_b m_t \sin 2\theta_b (-1)^{i+1} B_0(\mu^2, m_t^2, m_{\tilde{b}_i}^2) \right. \\ & \left. - |\mu| G_t^2 B_1(\mu^2, m_t^2, m_{\tilde{t}_i}^2) + |\mu| H_i^{tb} B_1(\mu^2, m_t^2, m_{\tilde{b}_i}^2) + (t \leftrightarrow b) \right] \\ & + \frac{\alpha}{\pi} |\mu| \left[B_0(\mu^2, \mu^2, 0) - B_0(\mu^2, \mu^2, M_Z^2) - \frac{1}{2} B_1(\mu^2, \mu^2, 0) + \frac{1}{2} B_1(\mu^2, \mu^2, M_Z^2) \right], \end{aligned} \quad (5)$$

where

$$\begin{aligned} H^{tb} &\equiv (G_t^2 \cos^2 \theta_b + G_b^2 \sin^2 \theta_b, G_t^2 \sin^2 \theta_b + G_b^2 \cos^2 \theta_b), \\ G_t &\equiv \sqrt{\frac{3G_F}{8\sqrt{2}\pi^2}} \frac{m_t}{\sin \beta}, \quad G_b \equiv \sqrt{\frac{3G_F}{8\sqrt{2}\pi^2}} \frac{m_b}{\cos \beta}. \end{aligned} \quad (6)$$

Explicit expressions for the functions B_0 and B_1 , defined as

$$B_n(p^2, m_1^2, m_2^2) = - \int_0^1 dx x^n \log[-p^2 x(1-x) + m_1^2(1-x) + m_2^2 x - i\epsilon] \quad (n = 0, 1), \quad (7)$$

can be found in ref. [9]. The stop mixing angle θ_t is defined such that $\tilde{t}_1 = \cos \theta_t \tilde{t}_L + \sin \theta_t \tilde{t}_R$ is the heavier mass eigenstate and $\tilde{t}_2 = -\sin \theta_t \tilde{t}_L + \cos \theta_t \tilde{t}_R$ is the lighter one; the sbottom mixing angle θ_b is defined analogously.

As expected, all terms in eqs. (4) and (5) vanish in the limit of exact electroweak symmetry. In this limit, all higgsino mass terms other than μ are forbidden. Notice that eq. (4) also vanishes when the stop mixing angle is zero (neglecting the term proportional to m_b), in spite of the presence of the electroweak breaking top-quark mass m_t . This happens because, in order to generate a non-vanishing $\delta\Delta_0$, one has to break an R -symmetry under

⁴The overall sign of $\delta\Delta_0$ is chosen by assuming that the lightest neutralino is determined by the tree-level relation in eq. (1).

which the Higgs superfields H_1 and H_2 and the top quark superfields Q_L, \bar{U}_R carry charges $R = \{2, 0, 0, 2\}$, respectively. This R -symmetry, although not broken by the top mass, is violated by stop left–right mixing terms.

The dominant contribution in eq. (4) comes from the top–stop loops and it is approximately given by

$$\delta\Delta_0 \simeq 2G_t^2 m_t \sin 2\theta_t \log \left[\frac{\max(m_t^2, m_{\tilde{t}_2}^2)}{m_{\tilde{t}_1}^2} \right], \quad (8)$$

for large mass splitting ($m_{\tilde{t}_1}^2 \gg m_{\tilde{t}_2}^2$). For maximal stop left–right mixing ($\sin 2\theta_t \simeq 1$) eq. (8) predicts $|\delta\Delta_0| \sim 12 \text{ GeV} \times \left(\frac{m_t}{180 \text{ GeV}}\right)^3 \left(\frac{1}{\sin^2 \beta}\right)$ for a stop mass $m_{\tilde{t}_1} \sim 1 \text{ TeV}$. Under these conditions, the one-loop corrections to Δ_0 can easily be of the order of the tree-level value. Notice, however, that these conditions on the stop mass parameters are not the ones that maximize the supersymmetric corrections to R_b [6]. The sign of the corrections depends on the sign of $\sin 2\theta_t$, which in turn is proportional to the unknown value of the stop left–right mixing. Therefore $\delta\Delta_0$ can either enhance or suppress the tree-level result.

The contributions proportional to $|\mu|$ in eq. (5) never amount to more than a few GeV for any value of $|\mu|$ relevant to LEP2 searches. Therefore when $\tan \beta$ is not too large, eq. (5) approximately predicts:

$$\delta\Delta_+ \simeq \frac{\delta\Delta_0}{2}, \quad (9)$$

which mimics the tree-level relation of eq. (3). On the other hand, for large $\tan \beta$, the second term in eq. (5) can become important and may destroy the correlation between $\delta\Delta_0$ and $\delta\Delta_+$ given by eq. (9).

Since the loop corrections to Δ_0 and Δ_+ are coming from electroweak breaking effects, it is necessary to check whether the same choice of parameters also leads to unacceptably large corrections to the electroweak observables at LEP1. We have verified that there are no large effects for either ϵ_2 or ϵ_3 [10] (or equivalently U or S [11]). For instance, taking maximal mass splittings $m_{\tilde{t}_1}^2 \gg m_{\tilde{t}_2}^2 > M_Z^2$ and $m_{\tilde{b}_1}^2 \gg m_{\tilde{b}_2}^2 > M_Z^2$, we find [12]

$$\begin{aligned} S = \frac{4 \sin^2 \theta_W}{\alpha} \epsilon_3 = \frac{1}{12\pi} & \left[\log \frac{m_{\tilde{b}_1}^2}{m_{\tilde{t}_1}^2} + \sin^2 \theta_t (4 - 3 \sin^2 \theta_t) \log \frac{m_{\tilde{t}_1}^2}{m_{\tilde{t}_2}^2} \right. \\ & \left. + \sin^2 \theta_b (2 - 3 \sin^2 \theta_b) \log \frac{m_{\tilde{b}_1}^2}{m_{\tilde{b}_2}^2} - \frac{5}{4} (\sin^2 2\theta_t + \sin^2 2\theta_b) \right], \end{aligned} \quad (10)$$

which is smaller than about 0.1, even for maximal squark left–right mixing and mass splittings as large as $m_{\tilde{t}_1}/m_{\tilde{t}_2} \sim m_{\tilde{b}_1}/m_{\tilde{b}_2} \sim 10$.

More important is the constraint coming from the ρ parameter. The contribution from stop and sbottom loops gives [13]

$$\begin{aligned} \Delta\rho = \frac{3G_F}{4\sqrt{2}\pi^2} & \left\{ \cos^2\theta_t \left[\cos^2\theta_b f(m_{\tilde{t}_1}^2, m_{\tilde{b}_1}^2) + \sin^2\theta_b f(m_{\tilde{t}_1}^2, m_{\tilde{b}_2}^2) \right] \right. \\ & + \sin^2\theta_t \left[\cos^2\theta_b f(m_{\tilde{t}_2}^2, m_{\tilde{b}_1}^2) + \sin^2\theta_b f(m_{\tilde{t}_2}^2, m_{\tilde{b}_2}^2) \right] \\ & \left. - \cos^2\theta_t \sin^2\theta_t f(m_{\tilde{t}_1}^2, m_{\tilde{t}_2}^2) - \cos^2\theta_b \sin^2\theta_b f(m_{\tilde{b}_1}^2, m_{\tilde{b}_2}^2) \right\}, \end{aligned} \quad (11)$$

where

$$f(x, y) = \frac{xy}{x-y} \log \frac{y}{x} + \frac{x+y}{2}. \quad (12)$$

In order to study the predictions on $\delta\Delta_+$ and $\delta\Delta_0$ compatible with the present constraint on the ρ parameter, we first define the stop and sbottom squark mass matrices as

$$\begin{aligned} m_{\tilde{t}}^2 &= \begin{pmatrix} m_{\tilde{Q}_L}^2 + m_t^2 + (\frac{1}{2} - \frac{2}{3}\sin^2\theta_W)\cos 2\beta M_Z^2 & m_t(A_t - \mu \cot\beta) \\ m_t(A_t - \mu \cot\beta) & m_{\tilde{t}_R}^2 + m_t^2 + \frac{2}{3}\sin^2\theta_W \cos 2\beta M_Z^2 \end{pmatrix}, \\ m_{\tilde{b}}^2 &= \begin{pmatrix} m_{\tilde{Q}_L}^2 + m_b^2 - (\frac{1}{2} - \frac{1}{3}\sin^2\theta_W)\cos 2\beta M_Z^2 & m_b(A_b - \mu \tan\beta) \\ m_b(A_b - \mu \tan\beta) & m_{\tilde{b}_R}^2 + m_b^2 - \frac{1}{3}\sin^2\theta_W \cos 2\beta M_Z^2 \end{pmatrix}. \end{aligned} \quad (13)$$

Since we do not want to rely on specific model-dependent assumptions, we will treat the supersymmetry-breaking parameters $m_{\tilde{Q}_L}^2$, $m_{\tilde{t}_R}^2$, $m_{\tilde{b}_R}^2$, A_t and A_b as free variables. The result of varying these five supersymmetry-breaking parameters (and the sign of μ) compatibly with the constraints $\Delta\rho < 1 \times 10^{-3}$ or 3×10^{-3} is shown in fig. 1. We have chosen $m_t = 180$ GeV and $|\mu| = 80$ GeV, but the dependence on $|\mu|$ is insignificant within the range of interest for LEP2. Figure 1a corresponds to the case of minimal $\tan\beta$ consistent with perturbativity up to the GUT scale, $\sin\beta \simeq m_t/(195 \text{ GeV})$. Figure 1b corresponds to the case of maximal $\tan\beta$, $\tan\beta \simeq m_t/m_b$. The stop left–right mixing parameter A_t plays a crucial role, since $\delta\Delta_0$ tends to zero in the limit $\sin 2\theta_t \rightarrow 0$. The largest effects on $\delta\Delta_0$ are obtained for maximal stop mixing and mass splitting, $m_{\tilde{Q}_L}^2 \sim m_{\tilde{t}_R}^2 \sim A_t m_t$. Figure 1 shows the comparison between the region of $\delta\Delta_+ - \delta\Delta_0$ values which can be obtained by requiring $|\bar{A}_t| \equiv |2A_t/(m_{\tilde{Q}_L} + m_{\tilde{t}_R})| < 1$ and that where $|\bar{A}_t| < 3$. For $|\bar{A}_t| < 1$, the regions in the $\delta\Delta_+ - \delta\Delta_0$ space where $\Delta\rho < 1 \times 10^{-3}$ and $\Delta\rho < 3 \times 10^{-3}$ are about the same.

Although constraints from the ρ parameter reduce the maximum values of $\delta\Delta_0$ with respect to the estimate in eq. (8), the effect of the one-loop corrections to Δ_0 can still be sizeable, possibly of the order of the tree-level contributions. The relation between $\delta\Delta_0$ and $\delta\Delta_+$ in eq. (9) is a good approximation for the small $\tan\beta$ (see fig. 1a), but deviations can appear for very large values of $\tan\beta$ (see fig. 1b).

Finally, we want to show how the values of Δ_+ and Δ_0 in eq. (1) can be modified when the μ parameter is allowed to be complex. The electric dipole moments of the neutron and electron put severe constraints on the phase of μ [14]. Nevertheless these constraints can be relaxed if the masses of the first generation of squarks and sleptons or the gaugino masses are much larger than μ . For example, if the first generation of squarks and sleptons masses are larger than ~ 1 TeV, no limits can be placed on the phase of μ . Such large masses are allowed in scenarios with non-universal soft masses without problems of fine tuning [15]. For a complex μ parameter, $\mu = |\mu|e^{i\varphi}$, and M real and positive, we find:

$$\begin{aligned}\Delta_0 &= \left| 2ra \frac{M_W^2}{M} + \frac{2 \sin 2\beta \cos \varphi}{r|\mu|} \left[b|\mu|^2 \frac{M_W^2}{M^2} + (r^2 - 1)a^2 \frac{M_W^4}{M^2} \right] \right| + \mathcal{O}(1/M^3), \\ \Delta_+ &= \frac{\Delta_0}{2} + \sin 2\beta \cos \varphi (a - 1) \frac{M_W^2}{M} + |\mu|(b - 1) \frac{M_W^2}{M^2} \\ &\quad + \left[1 - r^2 + (2r^2 - 2 - \cos^2 2\beta)a^2 \right] \frac{M_W^4}{2|\mu|M^2} + \mathcal{O}(1/M^3),\end{aligned}\tag{14}$$

where

$$r = \sqrt{1 - \sin^2 2\beta \sin^2 \varphi}.\tag{15}$$

By comparing the leading $1/M$ terms in eq. (14), we find $\Delta_0 > \Delta_+ > 0$, for any value of φ and M'/M . If we assume the unification condition $M'/M = 5 \tan^2 \theta_W / 3$, then $\Delta_+/\Delta_0 < 5/8$ for any φ . Notice that the effect of a non-trivial phase is always to reduce the value of Δ_0 with respect to the cases $\varphi = 0, \pi$. However, for large $\tan\beta$, the dependence of Δ_+ and Δ_0 on the phase φ disappears. This happens because, in the limit $\langle H_1 \rangle \rightarrow 0$, one can rotate away the phase φ from the chargino and neutralino mass matrices. On the contrary, for small $\tan\beta$, the effect of φ on Δ_+ and Δ_0 can be significant. In the limit $\tan\beta \rightarrow 1$ and $\varphi \rightarrow \pi/2$, the leading $1/M$ contribution to eq. (14) vanishes and therefore Δ_+ and Δ_0 are drastically reduced. In this case, Δ_+ can even become negative and the chargino may be the lightest supersymmetric particle.

In conclusion, we have shown that in the higgsino region ($M \gg M_W$, $\mu \sim M_W$) the one-loop corrections to the mass splittings Δ_+ and Δ_0 can be of the same order of the tree-level values. The effect is significant only if the mixing angle and the mass splitting of the two stop states are large. Depending on the sign of the stop mixing angle, these corrections can decrease or increase the mass splittings Δ_+ and Δ_0 . Thus, radiative corrections can make charginos and neutralinos very degenerate in mass for values of M much smaller than previously thought. The opposite can also happen, and radiative corrections can allow the discovery of charginos and neutralinos at LEP2 in parameter regions predicted to be problematic by the tree-level relations. The values of Δ_+ and Δ_0 can also be reduced if μ has a large phase and $\tan \beta \sim 1$.

If $\tan \beta$ is not too large, the one-loop corrections to Δ_+ and Δ_0 are correlated, see eq. (9) and fig. 1a, and mimic the tree-level relation, see eq. (3). Therefore the conclusion of ref. [5] that neutralino search is an important experimental tool in the study of the higgsino region remains valid. However the relation between Δ_+ , Δ_0 and the gaugino masses can be lost. Finally, if $\tan \beta$ is extremely large, sbottom loop corrections can partially spoil the $\Delta_+ - \Delta_0$ correlation, and higgsino mass splittings more critically depend on the parameter choice.

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References

- [1] G. Altarelli et al., *Interim Report on Physics Motivations for an Energy Upgrade of LEP2*, CERN-TH/95-151, CERN-PPE/95-78; G. Altarelli et al., *Report of the Workshop on Physics at LEP2*, in preparation.
- [2] J.L. Feng, N. Polonsky and S. Thomas, preprint SLAC-PUB-95-7050 (1995).
- [3] D.R. Claes for the D0 Coll., *Proceedings of the 10th Topical Workshop on Proton-Antiproton Collider*, Batavia, IL (USA), 9-13 May 1995; L. Nodulman for the CDF Coll., *Proceedings of the International Europhysics Conference on High Energy Physics* (HEP 95), Brussels, Belgium, 27 July-2 August 1995.
- [4] C.H. Chen, M. Drees and J. Gunion, preprint UCD-95-43 (1995).
- [5] S. Ambrosanio, B. Mele, M. Carena and C.E.M. Wagner, preprint CERN-TH/95-286 (1995).
- [6] G. Altarelli, R. Barbieri and F. Caravaglios, Phys. Lett. **B314** (1993) 357; J.D. Wells, C. Kolda and G.L. Kane, Phys. Lett. **B338** (1994) 219; D. Garcia, A. Jimenez and J. Sola, Phys. Lett. **B347** (1995) 309, 321; E **B351** (1994) 278.
- [7] H.E. Haber and G.L. Kane, Phys. Rep. **117** (1985) 75.
- [8] D. Pierce and A. Papadopoulos, Phys. Rev. **D50** (1994) 565; Nucl. Phys. **B430** (1994) 278; A.B. Lahanas, K. Tamvakis and N.D. Tracas, Phys. Lett. **B324** (1994) 387.
- [9] G. Passarino and M. Veltman, Nucl. Phys. **B160** (1979) 151.
- [10] G. Altarelli and R. Barbieri, Phys. Lett. **B253** (1991) 161.
- [11] M.E. Peskin and T. Takeuchi, Phys. Rev. Lett. **65** (1990) 964.
- [12] M. Drees, K. Hagiwara and A. Yamada, Phys. Rev. **D45** (1992) 1725.
- [13] R. Barbieri and L. Maiani, Nucl. Phys. **B224** (1983) 32; C.S. Lim, T. Inami and N. Sakai, Phys. Rev. **D29** (1984) 1488; Z. Hioki, Prog. Theor. Phys. **73** (1985) 1283; J.A. Grifols and J. Sola, Nucl. Phys. **B253** (1985) 47.

- [14] W. Buchmüller and D. Wyler, Phys. Lett. **B121** (1983) 321; J. Polchinski and M.B. Wise, Phys. Lett. **B125** (1983) 393.
- [15] S. Dimopoulos and G.F. Giudice, Phys. Lett. **B357** (1995) 573; A. Pomarol and D. Tommasini, preprint CERN-TH/95-207 (1995) (hep-ph/9507462).

Figure Captions

Fig. 1: Region of values in the $\delta\Delta_+ - \delta\Delta_0$ plane obtained by varying the supersymmetry-breaking parameters defined in eq. (13), with $m_t = 180$ GeV, $|\mu| = 80$ GeV, $\tan\beta = 2.4$ (fig. 1a) and $\tan\beta \simeq m_t/m_b$ (fig. 1b). The different lines correspond to different constraints on $|\bar{A}_{t,b}|$ and $\Delta\rho$.

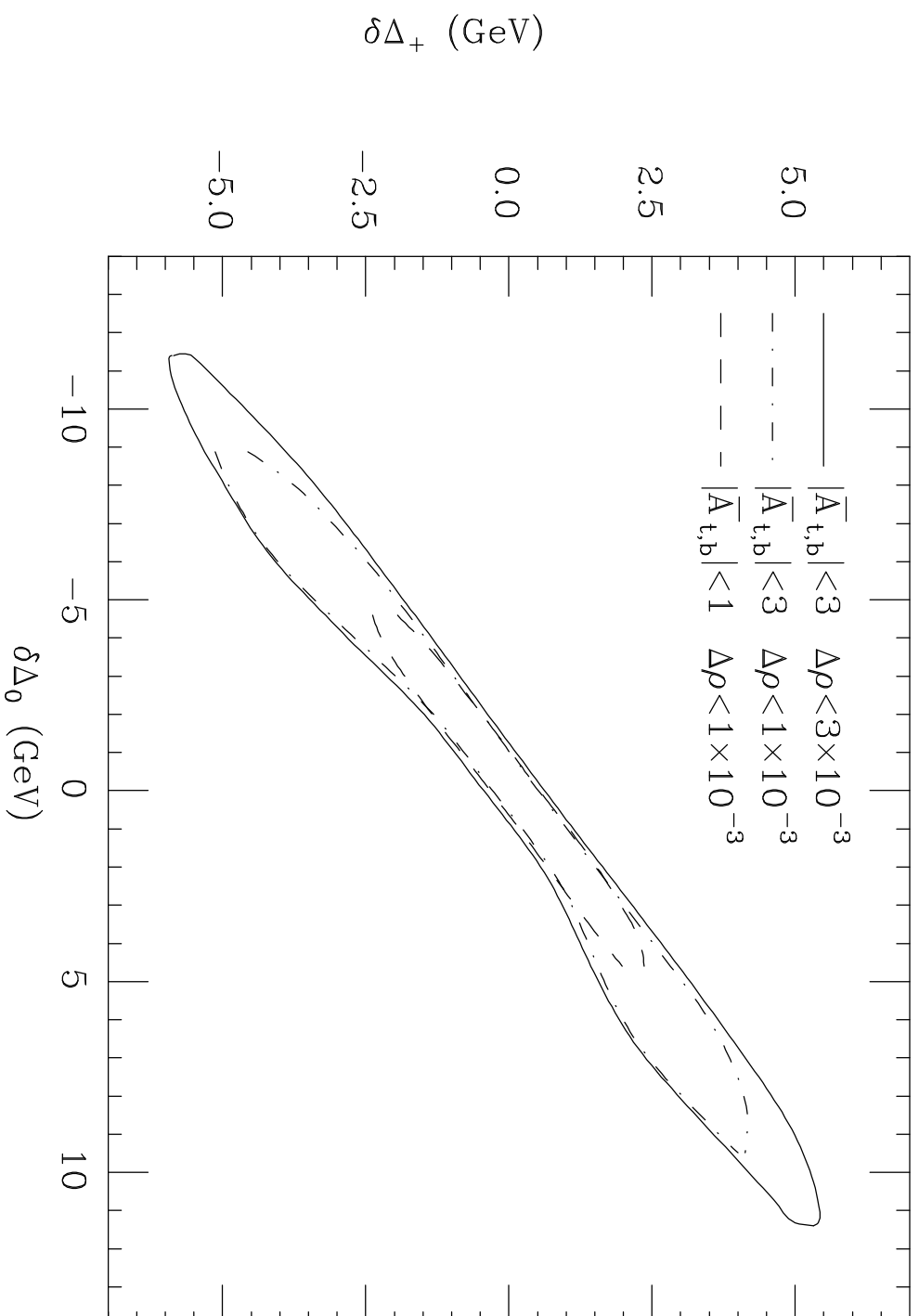


Fig. 1a

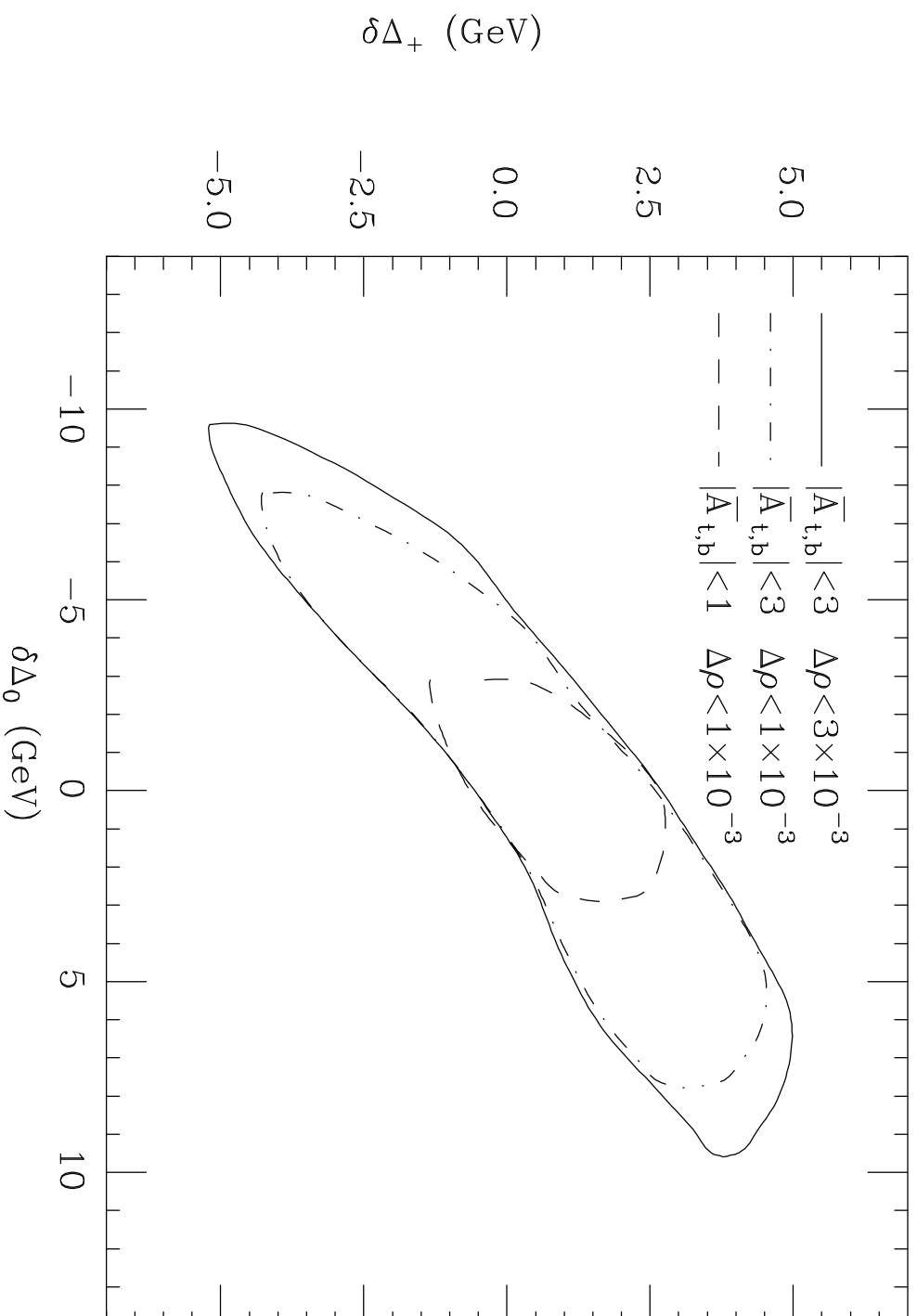


Fig. 1b